

A Supersymmetric Higgs Sector with Chiral D -terms

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ABSTRACT: Although supersymmetry remains the best candidate for solving the electroweak hierarchy problem, a supersymmetric Higgs boson near 125 GeV requires heavy scalars, highly-mixed stops, or non-minimal contributions to the Higgs potential. Extensions of the Standard Model (SM) gauge group provide an attractive means of raising the Higgs mass through non-decoupling D -term contributions to the Higgs quartic, but in most cases this correction is correlated with an enhanced coupling to bottom quarks and tau leptons that is disfavored by current fits to LHC Higgs data. In this work we demonstrate that the Higgs mass may be raised by non-decoupling D -terms without such enhanced couplings if the two supersymmetric Higgs doublets are “chiral”, i.e., charged under different gauge groups at high energies. In this case there is no direct correlation between the correction to the Higgs mass and its couplings to SM states, and in general the chiral correction to the Higgs potential undoes the MSSM preference for enhanced bottom couplings. This raises the prospects for discovering additional supersymmetric Higgs bosons consistent with the measured mass and couplings of the observed Higgs.

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1 Introduction

The recent discovery of a Higgs-like resonance [1, 2] and the first measurements of its branching ratios (BRs) provide a litmus test for models for physics beyond the Standard Model (SM). The massive vector decay modes $h \rightarrow WW^*$ and $h \rightarrow ZZ^*$ have been measured and agree well with the predictions for a SM-like Higgs. At the same time, hints of possible deviations from the SM predictions in the $h \rightarrow \gamma\gamma$ channel are fairly inconclusive and may well prove to be statistical fluctuations. The current fit of the Higgs BRs to the SM looks reasonable for the time being, strongly indicating that the Higgs is essentially SM-like and the deviations from the SM predictions are likely small.

The mass of the Higgs is somewhat perplexing from the point of view of new physics. While it is probably too light to be a “partially composite” Higgs, or to fit nicely within various generic technicolor-inspired models, it is also slightly too heavy to naturally fit the paradigm of minimal supersymmetry (SUSY). The minimal supersymmetric standard model (MSSM) at tree level predicts

$$m_h \leq m_Z |\cos(2\beta)| \tag{1.1}$$

due to the supersymmetric relation between the Higgs quartic and the electroweak gauge couplings. Radiative corrections dominated by stop-top loops can push the Higgs mass somewhat beyond this bound, but the corrections grow only logarithmically with the stop mass. A Higgs mass ~ 125 GeV requires stop masses of order 10 TeV, largely undermining the very motivation of SUSY as a solution to the hierarchy problem of the SM. Alternately, the Higgs mass may be raised by finite threshold corrections due to large stop mixing, but this requires a considerable numerical conspiracy between soft masses that is hard to achieve in many models of SUSY breaking. For a summary of the implications for the MSSM, see e.g. [3–11].

Supersymmetric naturalness – and the prospects for discovering additional states at the LHC – may be preserved in light of the observed Higgs mass if there are extra contributions to the Higgs quartic coupling above and beyond those of the MSSM. An effective field theory approach to such models was developed in [12], and concrete examples for these models include, among other possibilities, additional $SU(2)$ triplets [13, 14]; $SU(2)$ doublets [15, 16]; or gauge singlets (see [17] for review) that couple to the MSSM Higgses; or some combination thereof [18]. While none of these explanations is firmly excluded, in general they modify the Higgs branching ratios compared to the SM. It is challenging to find a parameter space for these models in which the lightest CP-even scalar masquerades effectively as the SM Higgs without any noticeable deviation from the measured branching ratios.

An alternate possibility is to enhance the SUSY Higgs quartic through additional non-decoupling D -terms, as first introduced in [19]. Such D -terms require a gauge extension of the MSSM group slightly above the TeV scale. For example, consider an $SU(2) \times SU(2)$ gauge group broken to a diagonal subgroup that we identify as the $SU(2)_L$ weak gauge group of the SM. If the MSSM Higgs doublets are originally charged under one of the $SU(2)$ ’s at high energies, they will have a D -term potential coming from both the unbroken SM gauge group and from the heavy W', Z' gauge fields of the broken group. Such models were studied in detail in [19–21] and found to possess various attractive properties, including almost-vanishing contributions to electroweak precision measurements and preservation of gauge coupling unification. Unfortunately these models also change the Higgs branching ratios, particularly the partial width $\Gamma(h \rightarrow b\bar{b})$. As shown in [22], these models are already in some tension with experimental data, which has only grown stronger as Higgs couplings are measured more accurately. While the enhancement of $\Gamma(h \rightarrow b\bar{b})$ may be decoupled by raising the masses of other scalars in the Higgs sector, this erases the prospects for discovering additional Higgs scalars at the LHC. This tension suggests D -term contributions to the potential, if responsible for raising the Higgs mass, should take an unconventional form.

In this work we will focus on non-decoupling D -terms that come from a conceptually different ultraviolet source, in which the MSSM Higgs doublets are chiral at high energies. A model along these lines was proposed in [23], but the full implications of this scenario for the Higgs sector remain unexplored. In this work we examine the effects of the non-decoupling chiral D -terms on the Higgs sector. We show that, unlike in the “standard” scenario of non-decoupling D -terms, corrections from the extended gauge sector tend to *suppress* the $h \rightarrow b\bar{b}$ and $h \rightarrow \tau^+\tau^-$ rates relative to the MSSM expectation. Since these rates are enhanced in the MSSM relative to the SM in the regime of large $\tan\beta$ [16], the chiral non-decoupling terms partially undo the “damage” to the Higgs potential caused by the MSSM itself. We also emphasize that in spite of the naive expectation, these chiral Higgs D -term can provide a good fit to precision electroweak measurements. In fact, the situation in the low-energy effective theory is improved compared to the MSSM.

The chiral Higgs scenario possesses a handful of non-trivial features. First, this scenario provides a natural resolution to the $\mu/B\mu$ problem. The conventional μ term is forbidden by gauge invariance in the full model, and only emerges beneath the scale of symmetry

breaking to the SM gauge group. Second, the $B\mu$ term can typically be very small, favoring very large values for $\tan\beta$. In general, $\tan\beta$ can vary from ~ 50 all the way to $\mathcal{O}(10^4)$, where one-loop non-holomorphic couplings dominate down-type quark and lepton fermion masses [24, 25]. When $\tan\beta \sim 50$ and tree-level couplings dominate, the corrections to Higgs couplings are unambiguous. At larger values of $\tan\beta$ where loop-suppressed non-holomorphic couplings dominate, there may be additional corrections to Higgs couplings with ambiguous signs.

Our paper is organized as follows: In Sec. 2 we review the standard non-decoupling D -term scenario and discuss its problem with $b\bar{b}$ and $\tau^+\tau^-$ Higgs decay modes. In Sec. 3 we introduce chiral Higgs D -terms and show that in the regime of moderately large $\tan\beta$ the corrections to the Higgs couplings fit SM expectations *better than the MSSM*. We also illustrate that the regime of very large $\tan\beta$ regime, namely $\tan\beta \sim \mathcal{O}(10^3 \dots 10^4)$, is well motivated in this scenario. We show that the Higgs branching ratios in this case do not significantly deviate from the SM predictions and, unlike in the moderate $\tan\beta$ regime, the corrections to the SM rate do not have a well-defined sign, but rather depend on various phases between soft terms. Finally, in Sec. 4 we conclude and suggest several other interesting questions related to this unusual “chiral” extension of the MSSM.

2 Non-decoupling D -terms with vector Higgses

We begin by reviewing the implications of conventional non-decoupling D -terms wherein both MSSM Higgs doublets are charged under the same gauge group at high energies. Consider a gauge extension of the SM, such that the full gauge group above the scale ~ 10 TeV is $SU(3)_c \times SU(2)_1 \times SU(2)_2 \times U(1)_1 \times U(1)_2$. It is also possible to double the $SU(3)$ group, in which case there are implications for mass mixing among colored sfermions. In this work we will not consider the $SU(3) \times SU(3)$ moose in detail since it is irrelevant for the Higgs discussion, but it is worth keeping this possibility in mind.

In addition to MSSM fields, which may be charged under various choices of these groups, there are bi-fundamental link fields $\chi, \tilde{\chi}$. The pairs of $SU(2)$ and $U(1)$ gauge groups are broken to their diagonal subgroups by VEVs of the link fields, which are in vector-like representations of $SU(2)_1 \times U(1)_1 \times SU(2)_2 \times U(1)_2$ and singlets of the color group, namely

$$\chi \sim (2_{1/2}, 2_{-1/2}), \quad \tilde{\chi} \sim (2_{-1/2}, 2_{1/2}). \quad (2.1)$$

The low-energy theory consists of the massless gauge fields of $SU(3)_c \times SU(2)_L \times U(1)_Y$, along with various massive states at the scale of the link field vevs.

The breaking of two replicas of the electroweak gauge group to the diagonal subgroup can be easily achieved, and may be largely supersymmetric. For example, consider the superpotential

$$W = \kappa \mathcal{A}(\chi\tilde{\chi} - f^2) \quad (2.2)$$

where \mathcal{A} is a gauge singlet dynamical field and κ is a dimensionless Yukawa coupling. This structure is particularly motivated by strong SUSY dynamics [26]. The fields χ , $\tilde{\chi}$ and \mathcal{A} also obtain soft masses from SUSY breaking. Provided these soft masses are $\lesssim f$, the

superpotential (2.2) triggers VEVs for χ , $\tilde{\chi}$ and breaks the extended gauge group down to $SU(2) \times U(1)$, which we identify with the electroweak group at low energies.¹

Although the low-energy gauge theory is that of the MSSM, novel phenomenology arises when MSSM fields are divided non-trivially among the gauge groups in the ultraviolet. Models along these lines [23, 27, 28] can at least partially explain the flavor puzzle of the SM, are attractive from the perspective of FCNCs [29], and lead to a “natural” SUSY spectrum [30, 31], making it non-trivial for LHC discovery and opening a window for SUSY particles below the TeV scale [32–34].

Whatever the configuration of matter fields, there is a potential correction to D -terms of states charged under the electroweak gauge group at low energies due to the massive gauge bosons. Unsurprisingly, this correction is proportional to supersymmetry breaking; if the higgsing of gauge groups is perfectly supersymmetric, the low-energy D -terms are those of the MSSM in accordance with supersymmetric decoupling theorems. However, in the presence of SUSY-breaking soft terms for the link fields, the integrating-out is no longer supersymmetric and the following hard SUSY breaking terms are induced in the effective potential:

$$V \supset \frac{g^2 \Delta}{8} \left(H_u^\dagger \sigma^a H_u + H_d^\dagger \sigma^a H_d \right)^2 + \frac{g'^2 \Delta'}{8} \left(H_u^\dagger H_u - H_d^\dagger H_d \right)^2, \quad (2.3)$$

where we have defined

$$\Delta \equiv \left(\frac{g_1}{g_2} \right)^2 \frac{M_s^2}{M_V^2 + M_s^2} \quad (2.4)$$

and a similar definition for Δ' , with substitutions $g \rightarrow g'$ and $M_V \rightarrow M'_V$. Here M_s stands for the soft mass of the sgoldstone mode of the link fields and M_V , M'_V are the masses of heavy W' and B' gauge bosons. As expected, these corrections are proportional to SUSY breaking and vanish in the SUSY limit.

This is a relatively clean and safe way to enhance the quartic coupling of the Higgs fields, which in turn raises the Higgs mass above the MSSM bound. In this case we obtain a new tree-level bound on the Higgs mass,

$$m_{h_{tree}}^2 \leq \left(m_Z^2 + \frac{g^2 \Delta + g'^2 \Delta'}{2} v^2 \right) \cos^2(2\beta) \quad (2.5)$$

where $v \approx 174$ GeV. In order to match the observed Higgs mass without requiring large radiative corrections, we require (in the large $\tan \beta$ limit)

$$\frac{g^2 \Delta + g'^2 \Delta'}{2} v^2 \sim \mathcal{O}(m_W^2). \quad (2.6)$$

¹Note that when SUSY is broken, in addition to quadratic soft terms we can also imagine a trilinear soft term $\mathcal{L} \supset a \mathcal{A} \chi \tilde{\chi} + \text{c.c.}$ This term would produce a tadpole for \mathcal{A} , shifting it from the origin and inducing non-vanishing F -terms for χ and $\tilde{\chi}$. In this particular analysis we have neglected this term because we are motivated by gauge-mediation like scenarios where the D -terms are usually small and therefore $a \ll \tilde{m}_\chi$. While there is nothing wrong with large F_χ in the Higgs analysis presented here, it can be more problematic when we try to embed it into a full theory that addresses the flavor puzzle. The situation can become even worse in Sec. 3, where a large F -term can induce an enormous $B\mu$ term and destabilize electroweak symmetry breaking. Therefore it is important to focus on UV completions in which this a -term is not too large.

Unfortunately, this beautiful mechanism of addressing the Higgs mass problem in supersymmetry has unintended consequences for the Higgs branching ratios. Let us focus on the large $\tan\beta$ limit, which is necessary to saturate the tree level bound on the Higgs mass from (2.5). A generic prediction of these non-decoupling D -terms at large $\tan\beta$ is an enhanced coupling of the Higgs to $b\bar{b}$ and $\tau^+\tau^-$ without significant modification of other important Higgs couplings such as VV and $t\bar{t}$. To show this, we very briefly review here relevant parts of a detailed analysis performed in [22, 35]. Similar conclusions were obtained in [16].

The MSSM Higgs sector is nothing but a two-Higgs-doublet model (2HDM) with a restricted set of parameters. The most generic renormalizable couplings of a 2HDM may be parameterized by the potential [36]:²

$$V = m_1^2|H_d|^2 + m_2^2|H_u|^2 + \frac{\lambda_1}{2}|H_d|^4 + \frac{\lambda_2}{2}|H_u|^4 + \lambda_3|H_u|^2|H_d|^2 + \lambda_4|H_d^\dagger H_u|^2 + \left(\frac{\lambda_5}{2}(H_u H_d)^2 + m_{12}^2 H_u H_d + \lambda_6|H_d|^2(H_u H_d) + \lambda_7|H_u|^2(H_u H_d) + \text{c.c.} \right). \quad (2.7)$$

In the decoupling limit with large $\tan\beta$, the heavy states are identified with H_d ; we can simply integrate this field out to obtain the couplings of fermions to the SM-like Higgs. If we neglect potential non-holomorphic couplings of the Higgs to $b\bar{b}$ and $\tau^+\tau^-$, we find the following expressions for the Higgs couplings to the down-type fermions:

$$r_{b/\tau} \equiv \frac{g_{hb\bar{b}/h\tau^+\tau^-}}{g_{hb\bar{b}/h\tau^+\tau^-}^{SM}} \approx \left(1 + \frac{m_h^2}{m_H^2}\right) \left(1 - \frac{2(\lambda_3 + \lambda_5)v^2}{m_H^2 - m_h^2}\right) + \dots \quad (2.8)$$

where the ellipses stand for terms that are usually negligible in SUSY models by virtue of being proportional to λ_7 or non-holomorphic couplings. The couplings to $b\bar{b}$ and $\tau^+\tau^-$ are the only Higgs couplings that are corrected at the leading order at tree level in the limit of large $\tan\beta$. Corrections to VV and $t\bar{t}$ couplings are suppressed by $1/\tan^2\beta$ and may be neglected.

In the MSSM with non-decoupling D -terms, the correction to the $b\bar{b}$ and $\tau^+\tau^-$ couplings takes the form

$$r_{b/\tau} \approx \left(1 + \frac{m_h^2}{m_H^2}\right) \left(1 + \frac{(g^2(1 + \Delta) + g'^2(1 + \Delta'))v^2}{2(m_H^2 - m_h^2)}\right). \quad (2.9)$$

Note that MSSM *without* any D -term corrections (i.e., $\Delta = \Delta' = 0$) already enhances the rates for $h \rightarrow b\bar{b}$ and $h \rightarrow \tau^+\tau^-$. The corrections to the Higgs quartic from non-decoupling D -terms further enhance these rates by an amount proportional to the tree-level contribution to the Higgs mass. This result is fairly generic, and although it is not yet excluded, it is somewhat in odds with the data. In particular, such enhancement is constrained by direct measurements of $h \rightarrow \tau^+\tau^-$ and $h \rightarrow b\bar{b}$, as well as indirect

²Note that this reference uses 2HDM conventions with equal hypercharge for the two doublets; we prefer opposite-hypercharge conventions for a SUSY analysis. One can simply switch from our conventions to conventions of [36] by identifying $\Phi_2 = H_u$, $\Phi_1 = i\sigma^2 H_d^*$.

limits on $h \rightarrow b\bar{b}$ coming from the measurement of Higgs decays to WW^* , ZZ^* and $\gamma\gamma$. This is illustrated in Figure 1, which shows the tree-level Higgs mass prediction and the enhancement of r_b relative to the Standard Model value.

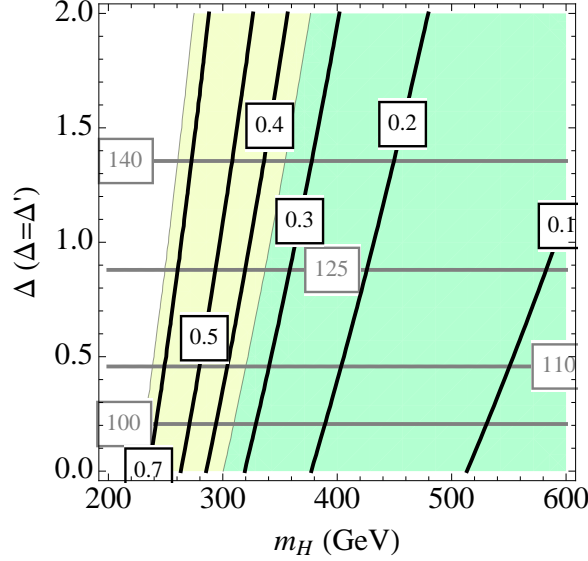


Figure 1. The tree-level prediction for the Higgs mass in GeV (gray lines) and the enhancement of the bottom coupling $r_b - 1$ above the Standard Model value (black lines) in the vector case as a function of the parameter Δ (taking $\Delta = \Delta'$) and the heavy CP-even Higgs mass m_H . Additional radiative contributions from stop loops further raise the Higgs mass, but are not required to accommodate a Higgs mass around 125 GeV. The green region contains the $+1\sigma$ fit to r_b using the most recent CMS results [37]; yellow contains the $+2\sigma$ fit. Points in the white region lie more than 2σ outside the CMS fit to r_b .

We can make this somewhat more quantitative by studying the current best fits to Higgs signals. Recently this has become a fairly precise procedure conducted directly by the collaborations, with full benefit of information about correlations between measurements in various channels. For simplicity we will focus our discussion on results from CMS [37], which to date present the most comprehensive information on Higgs coupling fits using $5 \oplus 12 \text{ fb}^{-1}$ of $7 \oplus 8 \text{ TeV}$ data.³ In a generic search for deviations in the six measured couplings (tree-level couplings to top quarks, bottom quarks, tau leptons, and vectors, plus effective couplings to gluons and photons), the 1D fit to the bottom quark coupling obtained by profiling the other five couplings yields $r_b \simeq 1.1^{+0.6}_{-0.5}$, while a similar fit to the tau coupling yields $r_\tau \simeq 0.95 \pm 0.4$. Although these coupling fits do not assume the particular relations among parameters present in the MSSM, they imply that current data is inconsistent with significant enhancement of couplings to down-type quarks. Similarly, in a search for deviations in the equivalence of couplings to down-type and up-type quarks (varying the coupling to vectors, the coupling to up-type quarks, and the coupling modifier ratio $\lambda_{du} \equiv r_b/r_t$), the 1D fit to λ_{du} profiling over the remaining couplings yields (when

³For a theory-level fit focusing on coupling to down-type quarks, see [16].

restricted to positive values of the couplings) $\lambda_{du} \simeq 1.05^{+0.35}_{-0.28}$, suggesting that the coupling to down-type quarks is not significantly enhanced over the coupling to up-type quarks.

Of course, the lack of evidence for enhanced bottom couplings may be reconciled with non-decoupling D -terms by raising the masses of the additional supersymmetric Higgs scalars. However this separation might come at the price of increasing the tree-level tuning in the Higgs potential unless a theoretically motivated mechanism is introduced to explain why $|m_{H_d}^2| \gg |m_{H_u}^2|$.⁴ From a pragmatic perspective, this also lessens the prospects for discovering additional Higgs scalars at the LHC if the couplings of the recently-discovered resonance prove largely SM-like.

3 Non-decoupling D -terms with chiral Higgses

The undesirable relation between vector-like D -terms and enhanced coupling to down-type quarks suggests a natural alternative to raise the Higgs mass: D -terms arising from Higgses charged under different gauge groups at high energies, which we will refer to as “chiral Higgses”. As before, consider a gauge group at high energies of the form $SU(3)_c \times SU(2)_1 \times SU(2)_2 \times U(1)_1 \times U(1)_2$, broken to the Standard Model around 10 TeV. However, instead of charging the Higgses under the same group, as in the previous section, consider charging the Higgses under *different groups*, namely

$$H_u \sim (2_{1/2}, 1_0), \quad H_d \sim (1_0, 2_{-1/2}) . \quad (3.1)$$

This seemingly innocent variation on the theme of non-decoupling D -terms has far-reaching consequences. Several immediate implications of this scenario are:

1. This model is necessarily a model of flavor, as gauge invariance necessarily forbids a subset of MSSM Yukawa couplings. These forbidden Yukawa couplings may be induced via irrelevant operators, which naturally provides a rudimentary flavor structure. The particular flavor structure depends on the distribution of MSSM matter fields between the two groups, which in turn is subject to constraints from anomaly cancelation.
2. A conventional MSSM μ -term is forbidden by gauge invariance. This is a desirable feature, since it automatically avoids the usual coincidence problem of the μ scale in the MSSM. Using the link fields defined in (2.1), the leading contribution to μ comes from the gauge invariant coupling

$$W = \lambda \tilde{\chi} H_u H_d \quad (3.2)$$

which leads to an effective μ -term when $\tilde{\chi}$ gets a VEV: $\mu_{eff} = \lambda \langle \tilde{\chi} \rangle$. This provides a reasonable μ term if $\lambda \lesssim 0.1$.

3. It is no longer transparent that custodial symmetry is preserved, and one must check that corrections to the ρ -parameter are under control.

⁴For particular examples of such models see e.g. [23, 38].

Two first two implications were addressed in [23], and we briefly revisit them here. The minimal way to assign the charges of the MSSM matter fields such that all the gauge anomalies cancel out is to take advantage of the fact that the gauge quantum numbers of the lepton doublet L and H_d in the MSSM are identical. Therefore we can simply cancel the gauge anomalies by charging a single L under $SU(2)_1 \times U(1)_1$, when all other fields of the same generation are charged under $SU(2)_2 \times U(1)_2$. In [23] we considered the second generation as a “split” one, while the entire third generation was charged under $SU(2)_1 \times U(1)_1$ and the entire first generation was charged under $SU(2)_2 \times U(1)_2$. This gave us a reasonable flavor structure in the limit $\tan \beta \sim \mathcal{O}(10^2 \dots 10^4)$, which we later justify in subsection 3.3. However, this charge assignment breaks full $SU(5)$ multiplets and therefore abandons hope of sensible gauge coupling unification. There is, however, a simple alternative that can preserve (accelerated) unification [39] and gives rise to an effective theory with “chiral” Higgses, obtained by “integrating in” the fields required to the generate Yukawa couplings that are otherwise forbidden by gauge invariance.

3.1 A simple model

The easiest way to obtain the chiral set of low-energy degrees of freedom with a sensible theory of flavor is to start in the UV with a pair of Higgs doublets at each $SU(2) \times U(1)$ – let’s call them A_u, A_d charged under $G_1 \equiv SU(2)_1 \times U(1)_1$, and B_u, B_d charged under $G_2 \equiv SU(2)_2 \times U(1)_2$; each pair has the quantum numbers of Higgs fields H_u, H_d and some GUT-scale mechanics has eliminated the triplet counterparts. We also need to assign SM fermions to either G_1 or G_2 . A particularly good choice is to charge the third generation under G_1 and the first two generations under G_2 ; this gives an order-unity top yukawa coupling and suppressed Yukawa couplings for the lighter generations, though typically the bottom yukawa coupling is suppressed. Another equally-valid possibility is to charge all three generations under G_1 ; this does not furnish a theory of flavor, but in this respect it is no worse off than the MSSM. We defer a detailed discussion of the various possibilities for later; the immediate consequences for the Higgs sector are insensitive to this choice.

The possible marginal or relevant superpotential terms allowed by the symmetries involving only the Higgs fields are:⁵

$$W \supset \mu_A A_u A_d + \mu_B B_u B_d + \lambda \tilde{\chi} A_u B_d + \lambda' \chi A_d B_u . \quad (3.3)$$

Now in the limit $\lambda' \langle \chi \rangle \gg \lambda \langle \tilde{\chi} \rangle, \mu_A, \mu_B$ we may simply integrate out A_d, B_u at tree level, leaving only a low-energy spectrum with $A_u \equiv H_u, B_d \equiv H_d$, and the low-energy superpotential

$$W = \left(\lambda \langle \tilde{\chi} \rangle - \frac{\mu_A \mu_B}{\lambda' \langle \chi \rangle} \right) H_u H_d + \dots \quad (3.4)$$

Down-type fermion bilinears charged under G_1 with tree-level couplings to A_d inherit couplings to H_d suppressed by $\mu_B / \lambda' \langle \chi \rangle$, while up-type fermion bilinears charged under G_2 with tree-level couplings to B_u inherit couplings to H_u suppressed by $\mu_A / \lambda' \langle \chi \rangle$. Thus all

⁵One might object that the presence of μ_A, μ_B reintroduces the μ -problem. However, this is far from a μ problem in the conventional sense; we require only that μ_A, μ_B are smaller than $\lambda' \langle \chi \rangle$, and may be arbitrarily small modulo the consequences for flavor.

Yukawa couplings and mixing angles may be generated without additional matter. This theory is manifestly anomaly-free above the scale $\lambda'\langle\chi\rangle$ and below the scale $\lambda\langle\tilde{\chi}\rangle$. In the intermediate regime, anomalies are compensated by Wess-Zumino terms in the usual way [40]. At high energies, if matter fields are distributed in unified multiplets then both groups possess the necessary matter content to preserve the low-energy supersymmetric prediction of (accelerated) gauge coupling unification [28]. Thus this simple model for chiral D -terms possesses all the features necessary for an anomaly-free theory of flavor, compatible with gauge-coupling unification. Many other variations are possible.

Now we may consider the non-decoupling D terms in the low-energy effective theory. The salient features are most apparent if we work in terms of the effective Kähler potential, which (for general abelian or non-abelian groups $G_1 \times G_2$) at low energies takes the form

$$K = H_u^\dagger e^{g_1 V_1} H_u + H_d^\dagger e^{g_2 V_2} H_d + \dots \quad (3.5)$$

in the basis where fundamental and anti-fundamental representations have the same generators. Now after higgsing, the massless vector field V_D and the massive vector field V_H are given by

$$V_D = \frac{g_1 V_2 + g_2 V_1}{\sqrt{g_1^2 + g_2^2}} \quad V_H = \frac{-g_1 V_1 + g_2 V_2}{\sqrt{g_1^2 + g_2^2}} \quad (3.6)$$

Then we can expand the Kahler potential to leading order in V_H to find

$$K \rightarrow H_u^\dagger e^{g_D V_D} H_u + H_d^\dagger e^{g_D V_D} H_d - \frac{g_D g_A}{g_B} H_u^\dagger V_H H_u + \frac{g_D g_B}{g_A} H_d^\dagger V_H H_d + \dots \quad (3.7)$$

Integrating out V_H with a spurion for the soft mass term leaves us with an effective action of the form

$$K_{eff} = H_u^\dagger e^{g_D V_D} H_u + H_d^\dagger e^{g_D V_D} H_d + g_D^2 \frac{M_s^2 \theta^4}{M_V^2 + M_s^2} \left(\frac{g_A}{g_B} H_u^\dagger T_D^a H_u - \frac{g_B}{g_A} H_d^\dagger T_D^a H_d \right)^2 + \dots \quad (3.8)$$

The relative sign between H_u and H_d in the correction term arises due to the relative sign in the coupling of the heavy vector superfield V_H to fields charged under G_1 and G_2 .

3.2 Spectrum and Higgs couplings in the moderate $\tan\beta$ regime

The contribution of (3.8) for the gauge groups $SU(2)_1 \times SU(2)_2 \times U(1)_1 \times U(1)_2$, with the charge assignments (3.1), leads to a potential of the form

$$\begin{aligned} \Delta V = & \frac{\xi}{2} \left(\omega |H_u^0|^2 + \omega |H_u^+|^2 + \frac{1}{\omega} |H_d^0|^2 + \frac{1}{\omega} |H_d^-|^2 \right)^2 + \\ & \frac{\xi'}{2} \left(\omega' |H_u^0|^2 + \omega' |H_u^+|^2 + \frac{1}{\omega'} |H_d^0|^2 + \frac{1}{\omega'} |H_d^-|^2 \right)^2 - \\ & 2\xi |H_u^+ (H_d^0)^* + H_u^0 (H_d^-)^*|^2 \end{aligned} \quad (3.9)$$

where for simplicity we have defined

$$\xi \equiv \frac{g^2}{4} \frac{M_s^2}{M_V^2 + M_s^2}, \quad \omega \equiv \frac{g_1}{g_2} \quad (3.10)$$

for the $SU(2)$ contributions, and analogously $\xi \rightarrow \xi', \omega \rightarrow \omega'$ for the hypercharge contribution. In the vector case we would have had, e.g., $\Delta = \frac{4}{g^2} \omega^2 \xi$, but in this case it is useful to separate the two factors. Note that in the limit of infinitely large supersymmetry breaking, the maximal values are $\xi \rightarrow \frac{g^2}{4}, \xi' \rightarrow \frac{g'^2}{4}$.

This new contribution modifies the usual tree-level relation between the β angle and the soft terms, which now takes the form

$$\sin 2\beta = \frac{2B\mu}{2|\mu|^2 + |m_u|^2 + |m_d|^2 + \xi v^2 f(\omega, \beta)} \quad (3.11)$$

where

$$f(\omega, \beta) \equiv \left(\omega + \frac{1}{\omega} \right) \left(\omega \sin^2 \beta + \frac{\cos^2 \beta}{\omega} \right). \quad (3.12)$$

The tree-level tadpole conditions also give a modified relation between m_Z and the the soft parameters, namely

$$m_Z^2 = \frac{|m_{H_d}^2 - m_{H_u}^2|}{\sqrt{1 - \sin^2(2\beta)}} - m_{H_u}^2 - m_{H_d}^2 - 2|\mu|^2 + \frac{2\xi v^2}{\cos(2\beta)} \left(\omega \sin^2 \beta + \frac{\cos^2 \beta}{\omega} \right) \left(\omega \sin^2 \beta - \frac{\cos^2 \beta}{\omega} \right) \quad (3.13)$$

noting that $\cos(2\beta) < 0$ for $m_{H_u}^2 < m_{H_d}^2$, as is the case here. At large $\tan \beta$ this becomes

$$m_Z^2 + 2(\xi\omega^2 + \xi'\omega'^2)v^2 \approx -2(|\mu|^2 + m_{H_u}^2) + \dots \quad (3.14)$$

which already indicates that the Higgs mass will accumulate a new positive contribution from the D -term corrections.

The pseudoscalar mass in this case is given by $m_A^2 = 2B\mu/\sin 2\beta$, as in the MSSM, but now the relation (3.11) implies a modification of the usual MSSM expression for m_A in terms of the Higgs soft masses:

$$m_A^2 = |\mu|^2 + m_u^2 + m_d^2 + \xi v^2 \left(\omega + \frac{1}{\omega} \right) \left(\omega \sin^2 \beta + \frac{\cos^2 \beta}{\omega} \right) + \xi' v^2 \left(\omega' + \frac{1}{\omega'} \right) \left(\omega' \sin^2 \beta + \frac{\cos^2 \beta}{\omega'} \right). \quad (3.15)$$

The charged Higgs mass is also deformed relative to the usual MSSM value by a negative contribution independent of g_1, g_2 :

$$m_{H^\pm}^2 = m_A^2 + m_W^2 - 2\xi v^2. \quad (3.16)$$

Note that this correction is independent of g_1, g_2 and brings the charged Higgses *closer* to the pseudoscalar in mass relative to the MSSM. If anything, this improves the fidelity of the

custodial $SU(2)$ triplet formed by H^\pm, A and reduces contributions to the ρ parameter in the low-energy effective theory. However, it does not reverse the ordering of the spectrum; in the limit of maximal supersymmetry breaking, it predicts exact degeneracy $m_{H^\pm} = m_A$. There are, of course, contributions to precision electroweak observables coming from heavy states [28, 41, 42], but these may be rendered sufficiently small for the mass scales of interest here.

The tree-level CP-even Higgs masses are given by

$$m_{H,h}^2 = \frac{1}{2} \left(m_A^2 + m_Z^2 + 2\xi v^2 \left(\omega^2 s_\beta^2 + \frac{c_\beta^2}{\omega^2} \right) + 2\xi' v^2 \left(\omega'^2 s_\beta^2 + \frac{c_\beta^2}{\omega'^2} \right) \pm \sqrt{\left(m_A^2 c_{2\beta} - m_Z^2 c_{2\beta} + 2\xi v^2 \left(\omega^2 s_\beta^2 + \frac{c_\beta^2}{\omega^2} \right) + 2\xi' v^2 \left(\omega'^2 s_\beta^2 + \frac{c_\beta^2}{\omega'^2} \right) \right)^2 + s_{2\beta}^2 (m_A^2 + m_Z^2 - 2\xi v^2)^2} \right), \quad (3.17)$$

which we may more usefully expand in the limit $\cos \beta \ll 1$ and $\xi \ll 1$. To the leading order in $1/\tan \beta$ and ξ , the tree-level masses are simply

$$m_h^2 = m_Z^2 + 2 \left(\xi \omega^2 + \xi' \omega'^2 \right) v^2 + \mathcal{O} \left(\frac{1}{\tan^2 \beta} \right) \quad (3.18)$$

$$m_H^2 = m_A^2 + \mathcal{O} \left(\frac{1}{\tan^2 \beta} \right). \quad (3.19)$$

Thus chiral D -terms persist in raising the tree-level prediction for the lightest CP-even Higgs mass, and may comfortably accommodate a Higgs around 125 GeV without reliance upon large radiative corrections.

In this respect, chiral D -terms do not differ significantly from conventional vector-like D -terms. However, they have strikingly different implications for the Higgs couplings. In particular, the correction from chiral D -terms to down-type quark and lepton couplings is now, in the limit of $m_h \ll m_H$ and $50 \gtrsim \tan \beta \gg 1$,

$$r_{b/\tau} = \left(1 + \frac{m_h^2}{m_H^2} \right) \left(1 + \frac{(g^2 + g'^2 - 4\xi - 4\xi')v^2}{2(m_H^2 - m_h^2)} \right). \quad (3.20)$$

Note the *negative* correction to the MSSM contribution. This result is independent of g_1, g_2 and, in the limit of maximal SUSY breaking, reduces precisely to $r_{b/\tau} \simeq 1 + m_h^2/m_H^2$, entirely undoing the effects of the MSSM quartic couplings. While the corrections from chiral D -terms cannot directly suppress the coupling to bottom quarks below the SM value, they essentially undo the MSSM preference for bottom enhancement. This raises the prospect that, if the couplings of the observed Higgs prove Standard Model-like, there may still be room for additional supersymmetric Higgs scalars at low energies.

3.3 Higgs couplings in the uplifted regime

Thus far we have focused on the situation at moderate $\tan \beta$, where tree-level Yukawa couplings dominate. However, as we have pointed out, it is likely that the $B\mu$ term vanishes at the scale $\langle \chi \rangle$ at which the extended gauge symmetry is higgsed. In this case the dominant contribution to the $B\mu$ term comes from the wino loop:

$$B\mu \approx -\frac{3\alpha_2^2}{2\pi} \mu m_{\tilde{W}} \ln \left(\frac{\langle \chi \rangle}{m_{\tilde{W}}} \right) + \dots \quad (3.21)$$

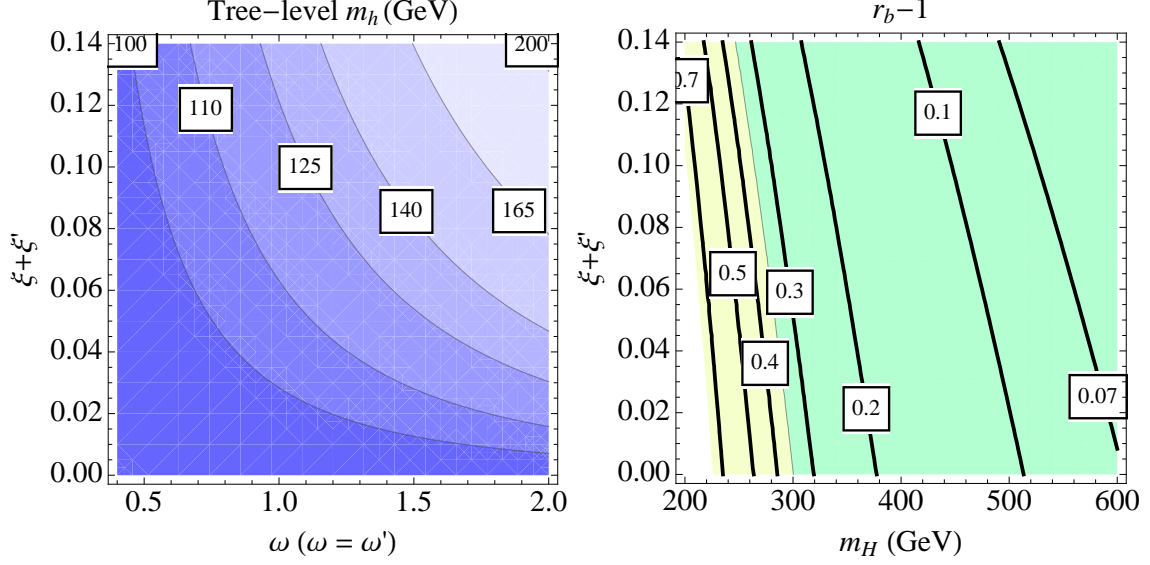


Figure 2. Left: The tree-level prediction for the Higgs mass as a function of the parameters ω (taking $\omega = \omega'$) and $\xi + \xi'$. Additional radiative contributions from stop loops further raise the Higgs mass, but are not required to accommodate a Higgs mass around 125 GeV. Right: The enhancement of the bottom coupling $r_b - 1$ above the Standard Model value as a function of the heavy CP-even Higgs mass m_H and the coupling $\xi + \xi'$. The green region contains the $+1\sigma$ fit to r_b using the most recent CMS results; yellow contains the $+2\sigma$ fit. Points in the white region lie more than 2σ outside the CMS fit to r_b . Note that the tree-level Higgs mass and couplings are parametrically liberated in this case.

This leads to the hierarchy $B\mu \ll \mu^2 \sim |m_{H_u}^2|$ and large values of $\tan\beta$. As we pointed out in [23], depending on the exact value of the wino mass, $\tan\beta$ in this case can vary between 40 and $\mathcal{O}(10^4)$. While $\tan\beta \sim \mathcal{O}(10)$ is definitely a viable possibility – in which case our analysis of couplings in the previous section is valid – the situation changes in the “uplifted” regime of $\tan\beta \gtrsim \mathcal{O}(100)$.

At such large values of $\tan\beta$, the vev of H_d is too small to provide masses for the down type quarks and the dominant contributions instead come from “wrong Yukawas”

$$\mathcal{L} \supset \hat{Y}_b H_u^\dagger Q b^c + \hat{Y}_\tau H_u^\dagger L \tau^c, \quad (3.22)$$

which are produced at the loop level. This scenario was analyzed in [24, 25] and it was found that in order to get sufficient quark masses one needs the “right” bottom and tau Yukawa couplings to be $\mathcal{O}(1)$, numerically exceeding the top Yukawa. This means that these Yukawa couplings reach Landau poles well below the GUT scale. On top of that, this particular part of parameter space is subject to severe flavor physics constraints, especially from $B \rightarrow \tau\nu$ and $B_s \rightarrow \mu^+\mu^-$. These constraints (as well as the issue of Landau poles for the Yukawa couplings) were carefully studied in [43]. Recently improved measurements of $B \rightarrow \tau\nu$ and $B_s \rightarrow \mu^+\mu^-$ [44] exacerbate constraints on SUSY at large $\tan\beta$. While detailed investigation of the numerical consequences of these new constraints is beyond the scope of this work, they merit mentioning before turning to a discussion of Higgs couplings.

Including such “wrong” Yukawa couplings, the relative enhancement of bottom couplings in a 2HDM (in an expansion in $1/\tan\beta$) takes the form [22]:

$$r_b = \frac{m_H^2}{m_H^2 - m_h^2} \left(1 - \frac{1}{1 + \hat{Y}_b \tan\beta} \left(\frac{2(\lambda_3 + \lambda_5)v^2}{m_H^2 - m_h^2} - \frac{2\lambda_7 v^2 \tan\beta}{m_H^2} + \frac{m_h^2}{m_H^2} \hat{Y}_b \tan\beta \right) \right), \quad (3.23)$$

where the couplings λ_3 , λ_7 and \hat{Y}_b are defined in Eqs. (2.7) and (3.22). Until now we have neglected the last two terms in this expression simply because they arise at one loop in SUSY models, and so are typically suppressed compared to the λ_3 term. However, when $\tan\beta \sim \mathcal{O}(100)$ we have $\lambda_3 \sim \lambda_7 \tan\beta \sim \hat{Y}_b \tan\beta$, and all the terms in this expression are equally important.⁶

Interestingly, the MSSM (as well as the chiral Higgs extension of the MSSM) does not have any firm prediction for the signs of λ_7 and \hat{Y}_b . Both these terms arise at the loop level. MSSM contributions to λ_7 have been calculated in [45] and (as expected) are determined by interference between three different diagrams; the sign of the result depends on the relative phase between the μ -term and the A -terms (in our language, it depends on the phase of λ in Eq. (3.2)). The situation is similar for the correction from \hat{Y}_b , which depends on the relative phases of soft parameters running in the loop. Therefore we conclude that the sign of new contributions to $\Gamma(h \rightarrow b\bar{b})$ and $\Gamma(h \rightarrow \tau^+\tau^-)$ is generally unfixed in the regime of very large $\tan\beta$, allowing for either positive and negative deviation depending on the size and importance of the λ_7 and \hat{Y}_b contributions.

Finally we comment briefly on the $h \rightarrow \gamma\gamma$ rate. Preliminary measurements of Higgs BRs by ATLAS and CMS suggested some deviations from the value predicted by the SM, although these deviations are not statistically significant and have shrunk with increasing data. However, it is still interesting to ask whether this particular model has any natural explanation for an enhanced diphoton rate. Since we are in the large $\tan\beta$ regime, it is natural to have a light and highly mixed stau. Note that the dominant contribution to stau mixing is proportional $Y_\tau \mu$, and Y_τ is relatively large already by $\tan\beta \sim 50$. As it was pointed out in [8] this can be a source of $h \rightarrow \gamma\gamma$ enhancement through new loop contributions to the $h\gamma\gamma$ coupling. Moreover, the enhancement need not be as large as in the MSSM, since contributions to the tree-level potential from chiral D -terms already undo much of the MSSM preference for bottom enhancement (and hence diphoton BR suppression).

3.4 Flavor and sflavor

Thus far we have remained entirely agnostic as to the distribution of MSSM matter superfields between gauge groups in the UV. This distribution may explain the broad features of Standard Model flavor and/or lead to a natural supersymmetric spectrum, depending on how SUSY breaking is mediated to both gauge groups. For completeness, we comment briefly on the possibilities:

⁶Note also that as expected all the new contributions are suppressed by m_H^2 , signaling that in general deviations are small, since one usually gets very large $\tan\beta$ for $|m_{H_d}^2| \gg |m_{H_u}^2|$.

- The “flavorless” scenario: All MSSM matter superfields are charged under G_1 . The Standard Model flavor hierarchy is explained through some external physics; the theory is no worse off than the MSSM from the perspective of flavor. The value of $\tan \beta$ depends on how SUSY breaking is communicated:
 1. SUSY breaking is mediated uniformly to both groups via gauge mediation with messengers charged under G_1 and G_2 . Sfermion masses have the usual gauge-mediated spectrum and $m_{H_d}^2$ is of the order of $m_{H_u}^2$, so that $\tan \beta$ is moderate.
 2. SUSY breaking is mediated preferentially to G_2 via gauge mediation with messengers charged under G_2 . Scalars charged under G_1 obtain mass via gaugino mediation, with a hierarchy relative to scalars charged under G_2 whose size depends on how much the gaugino masses are suppressed. Sfermion masses have a gaugino-mediated spectrum. However, $m_{H_d}^2$ may either be of the same order as $m_{H_u}^2$, in which case $\tan \beta$ is moderate, or it may be much larger, in which case $\tan \beta$ is large and the theory approaches the uplifted limit.
- The “flavorful” scenario: Third-generation MSSM matter superfields are charged under G_1 , while first- and second-generation matter superfields are charged under G_2 . This potentially can explain the broad features of the Standard Model flavor hierarchy without large hierarchies in dimensionless couplings, but only under certain circumstances depending on the communication of SUSY breaking:
 1. SUSY breaking is mediated uniformly to both groups via gauge mediation with messengers charged under G_1 and G_2 . Sfermion masses have approximately the usual gauge-mediated spectrum, but third-generation scalar masses may differ somewhat from first- and second-generation masses if the messenger numbers or couplings differ, leading to a mild “natural SUSY” hierarchy. Once again, $m_{H_d}^2$ is of the order of $m_{H_u}^2$, so $\tan \beta$ is moderate. However, in this case the tree-level flavor texture generated by marginal and irrelevant operators is poorly aligned with the observed pattern of fermion masses and mixings, so a suitable theory of flavor requires large hierarchies in dimensionless Yukawa couplings.
 2. SUSY breaking is mediated preferentially to G_2 via gauge mediation with messengers charged under G_2 . Scalars charged under G_1 obtain mass via gaugino mediation, with a hierarchy relative to scalars charged under G_2 whose size depends on how much the gaugino masses are suppressed. If gaugino masses are unsuppressed, there is a mild hierarchy between first/second- and third-generation scalar masses, $\tan \beta$ is moderate, and a suitable theory of flavor requires large hierarchies in dimensionless Yukawa couplings. If gaugino masses are significantly suppressed, there is a large hierarchy between first/second and third-generation scalar masses, $\tan \beta$ is large, and the uplifted limit is attained. Now the loop-induced Yukawa couplings provide a relatively good fit to the observed pattern of fermion masses and mixings, and no hierarchies in dimensionless Yukawa couplings are required.

Thus there are various viable configurations of MSSM matter superfields and messengers of gauge mediation, some of which may explain the broad features of the SM flavor hierarchy. There is typically a correlation between the separation of first/second- and third-generation soft masses and the value of $\tan \beta$; theories with moderate $\tan \beta$ typically have at most a mild, $\mathcal{O}(\text{few})$ hierarchy in the soft masses, while those in the uplifted regime with large $\tan \beta$ have a significant hierarchy between soft masses of the first two generations and the third generation.

4 Conclusions and Outlook

A Higgs at 125 GeV poses a sharp challenge for minimal supersymmetric extensions of the Standard Model and suggests either significant fine-tuning, large threshold corrections, or new contributions to the Higgs potential. Such new contributions, if present, may leave their imprint upon Higgs couplings. In the case of D -term contributions to the Higgs potential from extended gauge symmetries, the alteration of Higgs couplings typically enhances the partial width to bottom quarks relative to the Standard Model value. This is at variance with current fits to the Higgs couplings and suggests, if the Higgs ultimately proves Standard Model-like, that additional states in the Higgs sector must be significantly heavier than the observed resonance. In this work we have constructed a simple and explicit model for D -term corrections to the Higgs potential that instead favors Standard Model-like values for couplings to down-type quarks. This sensibly accommodates the observed Higgs mass and improves the prospects for achieving Standard Model-like decay rates (or perhaps an enhanced diphoton rate) even if the additional states in the Higgs sector are not decoupled in mass. These “chiral” D -terms automatically imply an approximate theory of flavor and a natural supersymmetric spectrum, as well as additional states that may be within distant reach at the LHC or a future hadron collider.

Acknowledgments

We are grateful to Alex Azatov, Kfir Blum, Spencer Chang, Jamison Galloway, Michele Redi, David Shih, and Scott Thomas for useful discussions. NC is supported in part by the DOE under grant DE-FG02-96ER40959, the NSF under grant PHY-0907744, and the Institute for Advanced Study. AK is supported by NSF grant No. PHY-0855591. We acknowledge the hospitality of the Aspen Center for Physics, supported by the National Science Foundation Grant No. PHY-1066293, where parts of this work were completed. AK also thanks the New High Energy Theory Center at Rutgers University for hospitality.

Note added: While this paper was being prepared for submission, [46] and [47] appeared. Both references discuss gauge extensions of the MSSM with vector-like charge assignments for the Higgs doublets; [46] emphasizes the implications for Higgs couplings from loop-level contributions, while [47] reiterates the tension between vector-like models and current fits to Higgs couplings.

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